

# Double polarization asymmetry as a possible filter for $\Theta^+$ 's parity

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We present an analysis of the Beam-Target double polarization asymmetry for the photoproduction of  $\Theta^+$  in  $\gamma n \rightarrow \Theta^+ K^-$ . We show that this quantity can serve as a filter for the determination of the  $\Theta^+$ 's spin-parity assignment near threshold. It is highly selective between  $1/2^+$  and  $1/2^-$  configurations due to dynamical reasons.

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The newly-discovered exotic  $\Theta^+$ , with strangeness  $S = +1$ , has stimulated a great number of theoretical discussions during the past few months. Although more experimental data are strongly recommended to confirm its properties, the significance of this discovery is that, for the first time, the existence of unconventional baryons is more than a purely theoretical prediction. So far, several experimental groups have seen this narrow state independently [1, 2, 3, 4, 5], while another state ( $\Xi^{--}$ ) was also reported by NA49 Collaboration recently [6].

Concerning the nature of the  $\Theta^+$ , its rather-low mass (1.535 GeV) and narrow width ( $< 25$  MeV) makes it an attractive candidate for the SU(3) Skyrme model predicted  $\bar{\mathbf{10}}$  multiplets [7, 8]. On the other hand, the existence of such an  $S = +1$  baryon does not rule out the conventional quark model, where baryons can be classified perfectly into  $\mathbf{8}$ ,  $\mathbf{10}$  and a singlet  $\mathbf{1}$ . Therefore, if an extra quark pair is present, e.g. a  $u\bar{s}$  within a neutron ( $udd$ ) or a  $d\bar{s}$  within a proton ( $uud$ ), a  $\Theta^+$  of pentaquark  $uudd\bar{s}$  can be constructed. However, such a treatment implies the extension of the SU(3) flavor space. For instance, multiplets of anti-decuplet  $\bar{\mathbf{10}}$ , 27-plet  $\mathbf{27}$ , and 35-plet  $\mathbf{35}$ , are possible, and will lead to different predictions in comparison with the Skyrme model.

The spin and parity of the exotic  $\Theta^+$  baryon turns out to be an essential issue for understanding the underlying dynamics. A series of theoretical studies exploring the underlying dynamics have been carried out, for instance, in the Skyrme model [9, 10, 11, 12, 13, 14, 15], quark models [16, 17, 18, 19, 20, 21, 22, 23, 24, 25, 26, 27, 28, 29], and other phenomenologies [30, 31], which rely crucially on the phenomenological assumptions for the  $\Theta^+$ 's quantum numbers. QCD sum rule studies [32, 33, 34, 35], and lattice QCD calculations [36, 37] are also reported. This situation suggests explicit experimental confirmation of the quantum numbers of the  $\Theta^+$  is not only important for establishing its status on a fundamental basis [38, 39], but also important for any progress in understanding its nature, and the existence of its other partners [40, 41].

A number of theoretical studies of the  $\Theta^+$  in meson photoproduction and meson-nucleon scattering were made recently concentrating on cross section predictions [42, 43, 44, 45, 46, 47]. However, due to the lack of knowledge about the underlying reaction mechanism, such studies of the reaction cross sections are strongly model-dependent. For instance, the total width of  $\Theta^+$  still has large uncertainties and could be much narrower [48, 49, 50, 51, 52, 53], and the role played by  $K^*$  exchange, as well as other  $s$ - and  $u$ -channel processes are unknown. Also, in a phenomenological approach the energy dependence of the couplings is generally introduced into the model via empirical form factors, which will bring further uncertainties. Taking this into account, there are advantages with polarization observables (e.g. ambiguities arising from the unknown form factors can be partially avoided). In association with the cross section studies, supplementary information about the  $\Theta^+$  can be obtained [54, 55, 56]. We also note that an interesting method to determine the  $\Theta^+$ 's parity was recently proposed in Ref. [57], and further detailed in Ref. [58].

In this letter, we will show that one of the beam-target (BT) double-polarization-asymmetry observables will be more selective to the  $\Theta^+$ 's parity. It is very likely that the BT asymmetry would serve as a parity filter for the  $\Theta^+$  near its production threshold. We will concentrate on the spin-parities of the  $\Theta^+$  of  $1/2^+$  and  $1/2^-$ , while some discussions will be devoted to the possible  $3/2$  partner.

For a positive parity  $\Theta^+$  ( $1/2^+$ ), a pseudovector effective Lagrangian is introduced for the  $\Theta NK$  coupling [54]. In  $\gamma n \rightarrow K^- \Theta^+$ , four transition amplitudes labelled by the Mandelstam variables will contribute in the Born approximation limit as shown by diagrams in Fig. 1: the contact term,  $t$ -channel kaon exchange,  $s$ -channel nucleon exchange and  $u$ -channel  $\Theta^+$  exchange. The transition can be expressed as:

$$\mathcal{M}_{fi} = M^c + M^t + M^s + M^u, \quad (1)$$

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where the four amplitudes are given by

$$\begin{aligned}
M^c &= ie_0 g_{\Theta NK} \bar{\Theta} \gamma_\mu \gamma_5 A^\mu N K, \\
M^t &= \frac{ie_0 g_{\Theta NK}}{t - M_K^2} \bar{\Theta} \gamma_\mu \gamma_5 (q - k)^\mu (2q - k)_\nu A^\nu N K, \\
M^s &= -g_{\Theta NK} \bar{\Theta} \gamma_\mu \gamma_5 \partial^\mu K \frac{[\gamma \cdot (k + P_i) + M_n]}{s - M_n^2} \left[ e_n \gamma_\alpha + \frac{i\kappa_n}{2M_n} \sigma_{\alpha\beta} k^\beta \right] A^\alpha N, \\
M^u &= -g_{\Theta NK} \bar{\Theta} \left[ e_\theta \gamma_\alpha + \frac{i\kappa_\theta}{2M_\Theta} \sigma_{\alpha\beta} k^\beta \right] A^\alpha \frac{[\gamma \cdot (P_f - k) + M_\Theta]}{u - M_\Theta^2} \gamma_\mu \gamma_5 \partial^\mu K N,
\end{aligned} \tag{2}$$

where  $e_0$  is the positive unit charge. The symbols,  $N$  and  $A$  denote the initial neutron and photon fields, while  $\bar{\Theta}$  and  $K$  denote the final state  $\Theta^+$  and kaon. In the  $s$ -channel the vector coupling vanishes since  $e_n = 0$ . We define the coupling constant  $g_{\Theta NK} \equiv g_A M_n / f_\theta$  with the axial vector coupling  $g_A = 5/3$ , while the decay constant is given by:

$$f_\theta = g_A \left( 1 - \frac{p_0}{E_n + M_n} \right) \left[ \frac{|\mathbf{p}'|^3 (E_n + M_n)}{4\pi M_\Theta \Gamma_{\Theta^+ \rightarrow K^+ n}} \right]^{1/2}, \tag{3}$$

where  $\mathbf{p}'$  and  $p_0$  are momentum and energy of the kaon in the  $\Theta^+$  rest frame, and  $E_n$  is the energy of the neutron. We adopt  $g_{\Theta NK} = 2.96$  corresponding to  $\Gamma_{\Theta^+ \rightarrow K^+ n} = 10$  MeV in the calculations. The  $\Theta^+$ 's magnetic moment  $\mu_\theta = 0.13(e_0/2M_\Theta)$  is estimated in Ref. [54] based on the the diquark model of Jaffe and Wilczek [17], which is consistent with the detailed model studies of Refs. [29, 35] and an estimate in Ref. [44].

The  $K^*$  exchange is also considered in this work. The effective Lagrangian for  $K^* K \gamma$  is given by

$$\mathcal{L}_{K^* K \gamma} = \frac{ie_0 g_{K^* K \gamma}}{M_K} \epsilon_{\alpha\beta\gamma\delta} \partial^\alpha A^\beta \partial^\gamma V^\delta K + \text{H.c.}, \tag{4}$$

where  $V^\delta$  denotes the  $K^*$  field;  $g_{K^* K \gamma} = 0.744$  is determined by the  $K^{*\pm}$  decay width  $\Gamma_{K^{*\pm} \rightarrow K^\pm \gamma} = 50$  keV [59].

The  $K^* N \Theta$  interaction is given by

$$\mathcal{L}_{\Theta NK^*} = g_{\Theta NK^*} \bar{\Theta} (\gamma_\mu + \frac{\kappa_\theta^*}{2M_\Theta} \sigma_{\mu\nu} \partial^\nu) V^\mu N + \text{H.c.}, \tag{5}$$

where  $g_{\Theta NK^*}$  and  $\kappa_\theta^*$  denote the vector and tensor couplings, respectively. So far, there is no experimental information about these two couplings. A reasonable assumption based on an analogue to vector meson exchange in pseudoscalar meson production is that  $|g_{\Theta NK^*}| = |g_{\Theta NK}|$ . For the tensor coupling, we assume  $|\kappa_\theta^*| = |\kappa_\rho| = 3.71$ , the same as  $\rho NN$  tensor coupling but with an arbitrary phase. Therefore, four sets of different phases are possible.

For  $\Theta^+$  of  $1/2^-$ , the effective Lagrangian for the  $\Theta^+ n K$  system conserves parity and is gauge invariant. This suggests that the electromagnetic interaction will not contribute to the contact term, which can be also seen in the leading term of the nonrelativistic expansion, where the derivative operator is absent. The Born approximation therefore includes three transitions (Fig. 2b, c and d): The invariant amplitude can be written as

$$\mathcal{M}_{fi} = M^t + M^s + M^u, \tag{6}$$

where the three transitions are given by

$$\begin{aligned}
M^t &= -\frac{e_0 g_{\Theta NK}}{t - M_K^2} \bar{\Theta} (2q - k)_\mu A^\mu N, \\
M^s &= g_{\Theta NK} \bar{\Theta} K \frac{[\gamma \cdot (k + P_i) + M_n]}{s - M_n^2} \left[ e_n \gamma_\mu + \frac{i\kappa_n}{2M_n} \sigma_{\mu\nu} k^\nu \right] A^\mu N, \\
M^u &= g_{\Theta NK} \bar{\Theta} \left[ e_\theta \gamma_\mu + \frac{i\kappa_\theta}{2M_\Theta} \sigma_{\mu\nu} k^\nu \right] A^\mu \frac{[\gamma \cdot (P_f - k) + M_\Theta]}{u - M_\Theta^2} N K,
\end{aligned} \tag{7}$$

where the coupling constant  $g_{\Theta NK} = [4\pi M_\Theta \Gamma_{\Theta^+ \rightarrow K^+ n} / |\mathbf{p}'| (E_n + M_n)]^{1/2}$  has a value of 0.61 corresponding to  $\Gamma_{\Theta^+ \rightarrow K^+ n} = 10$  MeV considered in this work.

If  $\Theta^+$  has spin-parity  $1/2^-$ , one may simply estimate its magnetic moment as the sum of  $(u\bar{s})$  and  $(udd)$  clusters, since the relative orbital angular momentum is zero. Assuming these two clusters are both color singlet, the total magnetic moment of this system can be written as

$$\mu_\theta = \left( \frac{2e_0}{6m_u} + \frac{e_0}{6m_s} \right) - \frac{e_0}{3m_u} = \frac{e_0}{6m_s}, \tag{8}$$

which leads to a small anomalous magnetic moment (since  $3m_s \simeq M_\Theta$ ). However, since such a simple picture may not be sufficient, we also include  $\kappa_\theta = \pm\kappa_p = \pm 1.79$  to make a sensitivity test.

We also include  $K^*$  exchange in the  $1/2^-$  production, in which the  $K^*N\Theta$  interaction is given by

$$\mathcal{L}_{\Theta NK^*} = g_{\Theta NK^*} \bar{\Theta} \gamma_5 (\gamma_\mu + \frac{\kappa_\theta^*}{2M_\Theta} \sigma_{\mu\nu} \partial^\nu) V^\mu N + \text{H.c.} \quad (9)$$

As in the  $1/2^+$  case, we assume  $|g_{\Theta NK^*}| = |g_{\Theta NK}| = 0.61$ . In the calculation we choose, somewhat arbitrarily, the anomalous magnetic moment  $|\kappa_\theta^*| = 0.37$ , which is one order of magnitude smaller than that in the production of  $1/2^+$ . Of course, we could have chosen it to be in the same ratio as that of  $g_{\Theta NK}$  in the two cases (about a factor of 5). However, the precise value of  $|\kappa_\theta^*|$  does not alter the overall conclusions of this paper regarding the qualitative features of the BT asymmetry near threshold.

In the helicity frame, the transition amplitude can be written as:

$$T_{\lambda_\theta, \lambda_\gamma \lambda_N} \equiv \langle \Theta^+, \lambda_\theta, \mathbf{P}_\theta; K^-, \lambda_0, \mathbf{q} | \hat{T} | n, \lambda_N, \mathbf{P}_i; \gamma, \lambda_\gamma, \mathbf{k} \rangle, \quad (10)$$

where  $\lambda_\gamma = \pm 1$ ,  $\lambda_N = \pm 1/2$ ,  $\lambda_0 = 0$ , and  $\lambda_\theta$  are helicities of photon, neutron,  $K^-$ , and  $\Theta^+$ , respectively. Following the convention of Ref. [60], we define the four independent helicity as:

$$\begin{aligned} H_1 &= T_{-\frac{1}{2}, 1-\frac{1}{2}}, \\ H_2 &= T_{-\frac{1}{2}, 1\frac{1}{2}}, \\ H_3 &= T_{\frac{1}{2}, 1-\frac{1}{2}}, \\ H_4 &= T_{\frac{1}{2}, 1\frac{1}{2}}. \end{aligned} \quad (11)$$

The BT polarization asymmetry is defined as the ratio between the polarized cross section and unpolarized one in terms of the total c.m. scattering angle  $\theta_{c.m.}$  (the angle between  $\mathbf{k}$  and  $\mathbf{q}$ ). Analytically, it can be related to the density matrix elements for  $\Theta^+ \rightarrow K^+ n$ . We are interested in the observable that has the photon circularly polarized along the photon moment direction  $\hat{z}$  and the neutron transversely polarized along the  $\hat{y}$ -axis perpendicular to the reaction plane. The  $\Theta^+$  decay density matrix element is defined as

$$\rho_{\lambda_\theta \lambda'_\theta}^{BT} = \frac{1}{2N} \sum_{\lambda_\gamma, \lambda_N} \lambda_\gamma T_{\lambda_\theta, \lambda_\gamma - \lambda_N} T_{\lambda'_\theta, \lambda_\gamma \lambda_N}^* \quad (12)$$

where,  $N = \frac{1}{2} \sum_{\lambda_\theta, \lambda_\gamma, \lambda_N} |T_{\lambda_\theta, \lambda_\gamma \lambda_N}|^2$  is the normalization factor. Because of parity conservation and the requirement of Hermiticity, only two elements are independent:

$$\begin{aligned} \rho_{\frac{1}{2}, \frac{1}{2}}^{BT} &= \rho_{-\frac{1}{2}, -\frac{1}{2}}^{BT}, \\ \rho_{-\frac{1}{2}, \frac{1}{2}}^{BT} &= -\rho_{\frac{1}{2}, -\frac{1}{2}}^{BT}, \quad \text{with} \quad \text{Re} \rho_{-\frac{1}{2}, \frac{1}{2}}^{BT} = \text{Re} \rho_{\frac{1}{2}, -\frac{1}{2}}^{BT} = 0. \end{aligned} \quad (13)$$

This leads to an expression for the BT asymmetry,

$$\mathcal{D}_{xz} = \frac{\rho_{\frac{1}{2}, \frac{1}{2}}^{BT}}{\rho_{\frac{1}{2}, \frac{1}{2}}^0} = \frac{1}{N} \text{Re} \{ H_1 H_2^* + H_3 H_4^* \}, \quad (14)$$

where  $\rho_{\frac{1}{2}, \frac{1}{2}}^0$  is the unpolarized density matrix element, and the subscript  $xz$  denotes the polarization direction of the the initial neutron target along  $x$ -axis in the production plane and the incident photon along the  $z$ -axis.

The expression of BT asymmetry for  $1/2^+$  and  $1/2^-$  is the same. But the underlying dynamics will be determined by the parities, and leads to different behaviors of  $\mathcal{D}_{xz}$ . For the production of  $1/2^+$  and  $1/2^-$  respectively, we will show that analytical features due to the dynamics arise from the BT asymmetry, and turn out to be quite different for the two parity cases.

For  $1/2^+$ , the transition amplitudes can be expressed in terms of the CGLN amplitudes [61]:

$$\langle \Theta^+, \lambda_\theta, \mathbf{P}_\theta; K^-, \lambda_0, \mathbf{q} | \hat{T} | n, \lambda_N, \mathbf{P}_i; \gamma, \lambda_\gamma, \mathbf{k} \rangle = \langle \lambda_\theta | \mathbf{J} \cdot \boldsymbol{\epsilon}_\gamma | \lambda_N \rangle, \quad (15)$$

where the operator  $\mathbf{J} \cdot \boldsymbol{\epsilon}_\gamma$  has a form:

$$\mathbf{J} \cdot \boldsymbol{\epsilon}_\gamma = i f_1 \boldsymbol{\sigma} \cdot \boldsymbol{\epsilon}_\gamma + f_2 \frac{1}{|\mathbf{q}| |\mathbf{k}|} \boldsymbol{\sigma} \cdot \mathbf{q} \boldsymbol{\sigma} \cdot (\mathbf{k} \times \boldsymbol{\epsilon}_\gamma) + i f_3 \frac{1}{|\mathbf{q}| |\mathbf{k}|} \boldsymbol{\sigma} \cdot \mathbf{k} \mathbf{q} \cdot \boldsymbol{\epsilon}_\gamma + i f_4 \frac{1}{|\mathbf{q}|^2} \boldsymbol{\sigma} \cdot \mathbf{q} \mathbf{q} \cdot \boldsymbol{\epsilon}_\gamma. \quad (16)$$

The coefficients  $f_{1,2,3,4}$  are functions of energies, momenta, and scattering angle  $\theta_{c.m.}$ , and contain information on dynamics. They provide an alternative expression for the BT asymmetry:

$$\mathcal{D}_{xz} = \sin \theta_{c.m.} \text{Re}\{f_1 f_3^* - f_2 f_4^* + \cos \theta_{c.m.} (f_1 f_4^* - f_2 f_3^*)\}. \quad (17)$$

Note that the above expression is the same as Eq. (B8) of Ref. [62].

Near threshold, useful analytical information can be obtained. It is well-established that the contact and  $t$ -channel terms are the leading contributions. In particular, the Kroll-Ruderman term ( $f_1$ ), is dominant in the transition. Compared to  $f_3$ , amplitude  $f_4$  is relatively suppressed by a further power on the (small) final-state three momentum  $|\mathbf{q}|$ . Amplitude  $f_2$  is also relatively suppressed in comparison with  $f_3$  as it comes from the magnetic transition. Therefore,  $f_1 f_4^*$  and  $f_2 f_3^*$  will both be relatively suppressed in comparison with  $f_1 f_3^*$  and further suppressed by  $\cos \theta_{c.m.}$  in the middle angles. Thus, the behaviour of  $\mathcal{D}_{xz}$  near threshold can be approximated by

$$\mathcal{D}_{xz} \simeq \sin \theta_{c.m.} \text{Re}\{f_1 f_3^*\}. \quad (18)$$

Since the CGLN amplitudes only depend weakly on the scattering angle  $\theta_{c.m.}$  (via the Mandelstam variables), this approximation implies a  $\sin \theta_{c.m.}$  behavior of  $\mathcal{D}_{xz}$ , and the sign of  $\mathcal{D}_{xz}$  is determined by the product.

In Fig. 2, as shown by the curves from full calculations at  $W = 2.1$  GeV, a clear  $\sin \theta_{c.m.}$  behaviour appears in the BT asymmetry. Although the sign change of the  $K^*$  exchange results in a quite significant change to the asymmetry values, the curves nevertheless confirm the dominance of the  $f_1 f_3^*$  term in Eq. (17). This result is by no means trivial. It suggests that although the  $K^*$  exchange might produce significantly different predictions for the cross sections [54], the BT asymmetry will tend to behave predominantly as  $\sin \theta_{c.m.}$  with its sign determined by  $f_1 f_3^*$ . Such a feature can be regarded as a signature of  $1/2^+$  spin-parity for the  $\Theta^+$ .

It is natural to expect that such a behavior should not hold at higher energies, where other mechanisms may contribute, and  $f_1$  and  $f_3$  will no longer be the leading terms. As shown by the asymmetries for different  $K^*$  exchange phases at  $W = 2.5$  GeV, structures deviating from  $\sin \theta_{c.m.}$  arise. In particular, comparing the dashed curve in Fig. 2 (b) and the dotted one in Fig. 2(d) with the solid curve for the exclusive Born terms, the structures reflect the strong interference from the  $K^*$  exchange via the  $\cos \theta_{c.m.} (f_1 f_4^* - f_2 f_3^*)$  term in Eq. (17).

Similar analysis can be applied to the production of  $1/2^-$ . In general, the transition amplitude can be arranged in a way similar to the CGLN amplitude:

$$\mathbf{J} \cdot \boldsymbol{\epsilon}_\gamma = iC_1 \frac{1}{|\mathbf{k}|} \boldsymbol{\sigma} \cdot (\boldsymbol{\epsilon}_\gamma \times \mathbf{k}) + C_2 \frac{1}{|\mathbf{q}|} \boldsymbol{\sigma} \cdot \mathbf{q} \boldsymbol{\sigma} \cdot \boldsymbol{\epsilon}_\gamma + iC_3 \frac{1}{|\mathbf{q}||\mathbf{k}|^2} \boldsymbol{\sigma} \cdot \mathbf{k} \mathbf{q} \cdot (\boldsymbol{\epsilon}_\gamma \times \mathbf{k}) + iC_4 \frac{1}{|\mathbf{q}|^2 |\mathbf{k}|} \boldsymbol{\sigma} \cdot \mathbf{q} \mathbf{q} \cdot (\boldsymbol{\epsilon}_\gamma \times \mathbf{k}), \quad (19)$$

where coefficients  $C_{1,2,3,4}$  are functions of energies, momenta and scattering angle, and contain dynamical information on the transitions. Restricted to the kinematics near threshold, a term proportional to  $\mathbf{q} \cdot \boldsymbol{\epsilon}_\gamma \boldsymbol{\sigma} \cdot (\mathbf{k} \times \mathbf{q})$  in the  $u$ -channel is neglected since it comes from a higher order contribution. An advantage of this formulation is that one can express the above in parallel to the CGLN amplitudes. Since  $(\boldsymbol{\epsilon}_\gamma \times \hat{\mathbf{k}}) = i\lambda_\gamma \boldsymbol{\epsilon}_\gamma$ , one can replace vector  $(\boldsymbol{\epsilon}_\gamma \times \hat{\mathbf{k}})$  with  $i\lambda_\gamma \boldsymbol{\epsilon}_\gamma$ , and rewrite the operator as

$$\mathbf{J} \cdot \boldsymbol{\epsilon}_\gamma = i\lambda_\gamma \left[ iC_1 \boldsymbol{\sigma} \cdot \boldsymbol{\epsilon}_\gamma + C_2 \frac{1}{|\mathbf{q}||\mathbf{k}|} \boldsymbol{\sigma} \cdot \mathbf{q} \boldsymbol{\sigma} \cdot (\mathbf{k} \times \boldsymbol{\epsilon}_\gamma) + iC_3 \frac{1}{|\mathbf{q}||\mathbf{k}|} \boldsymbol{\sigma} \cdot \mathbf{k} \mathbf{q} \cdot \boldsymbol{\epsilon}_\gamma + iC_4 \frac{1}{|\mathbf{q}|^2} \boldsymbol{\sigma} \cdot \mathbf{q} \mathbf{q} \cdot \boldsymbol{\epsilon}_\gamma \right], \quad (20)$$

which has exactly the same form as Eq. 16 apart from an overall phase factor from the photon polarization  $i\lambda_\gamma$ . It also suggests that the BT asymmetry for  $1/2^-$  in terms of those coefficients has the same form as Eq. 17:

$$\mathcal{D}_{xz} = \sin \theta_{c.m.} \text{Re}\{C_1 C_3^* - C_2 C_4^* + \cos \theta_{c.m.} (C_1 C_4^* - C_2 C_3^*)\}, \quad (21)$$

which is indeed the case. Quite remarkably, the behaviour of  $\mathcal{D}_{xz}$  due to these two different parities now becomes more transparent since the role played by the dynamics has been isolated out.

Note that since  $\mathcal{D}_{xz}$  is roughly proportional to  $\sin \theta_{c.m.}$ , it will vanish at  $\theta_{c.m.} = 0^\circ$  and  $180^\circ$ , and any sign change will reflect the competition among the  $C$  coefficients due to the dynamics.

The most important difference between the  $1/2^-$  and  $1/2^+$  cases is that the role of  $C_1$  may not be as significant as  $f_1$  in the production of  $1/2^+$ . In the Born approximation limit, the term of  $C_1$ , though dominant, comes from the  $s$ - and  $u$ -channel, which differs from the Kroll-Ruderman contribution from the contact term in  $1/2^+$  production. As a result, a relative suppression from the baryon propagator is expected when the energy increases. One thus may conjecture that other mechanisms, such as  $K^*$  exchange, may easily compete with the Born contribution and produce sign changes to the asymmetries above threshold.

However, as shown by the solid curve at  $W = 2.1$  GeV in Fig. 3, the BT asymmetry in the Born limit exhibits a  $\sin \theta_{c.m.}$  behavior, which suggests the dominance of either  $C_1 C_3^*$  or  $C_2 C_4^*$  near threshold, and the  $\text{Re}\{\cos \theta_{c.m.} (C_1 C_4^* -$

$C_2 C_3^*$ ) term should not be important. The inclusion of the  $K^*$  exchange certainly does not change this situation as shown by the dashed and dotted curves at  $W = 2.1$  GeV. In particular, an absolute sign difference, unlikely to be accidental, appears between these two parities, and needs to be understood.

First, let us try to understand the behavior of  $\mathcal{D}_{xz}$  in  $1/2^+$  production. As mentioned previously, the BT asymmetry for  $1/2^+$  is controlled by  $f_1 f_3^*$  in the Born terms. A detailed analysis gives the leading term of this combination:

$$f_1 f_3^* \simeq -e_0^2 g_{\Theta NK}^2 \frac{2}{u - M_\Theta^2} \mathcal{F}_c(k, q) \mathcal{F}_u(k, q), \quad (22)$$

where  $f_1$  is dominated by the Kroll-Ruderman term and  $f_3$  by the electric coupling in the  $u$ -channel. It is worth noting that the  $s$ -channel will not contribute to  $f_3$ , and the magnetic coupling in the  $u$ -channel is relatively suppressed. These features fix the sign of  $\mathcal{D}_{xz}$ , and underline the dominance of  $f_1 f_3^*$  a dynamical consequence.  $\mathcal{F}_c(k, q)$  and  $\mathcal{F}_u(k, q)$  are form factors for the contact and  $u$ -channel, and are treated the same in this approach. This assumption may bring in uncertainties at high energies, but should be a reasonable one near threshold.

For the production of  $1/2^-$ , we found that the  $C_1 C_3^*$  term plays a dominant role in the asymmetry. A detailed analysis shows that the dominant contribution to  $C_1$  is from the  $s$ -channel, while the contribution from the  $u$ -channel is strongly suppressed by the limited kinematics. A very important property arising from the  $C_3$  term is that it will be dominated by the  $t$ -channel kaon exchange via the decomposition  $\mathbf{q} \cdot \boldsymbol{\epsilon}_\gamma = \boldsymbol{\sigma} \cdot \mathbf{q} \boldsymbol{\sigma} \cdot \boldsymbol{\epsilon}_\gamma + i \boldsymbol{\sigma} \cdot (\boldsymbol{\epsilon}_\gamma \times \mathbf{k}) \mathbf{q} \cdot \mathbf{k} - i \boldsymbol{\sigma} \cdot \mathbf{k} \mathbf{q} \cdot (\boldsymbol{\epsilon}_\gamma \times \mathbf{k})$ . Therefore, we have

$$C_1 C_3^* \simeq -e_0^2 g_{\Theta NK}^2 \frac{\kappa_n}{2M_n} \frac{2|\mathbf{k}||\mathbf{q}|}{(W - M_n)(t - M_K^2)} \mathcal{F}_t(k, q) \mathcal{F}_s(k, q), \quad (23)$$

which will be negative since  $\kappa_n = -1.91$ . In comparison with Eq. 22, it gives a dynamical reason for the sign difference between the two parities. Also, note that  $\mathcal{F}_t(k, q)$  and  $\mathcal{F}_s(k, q)$  are form factors for the  $t$ - and  $s$ -channels, which is a further indication of the very different characteristic dynamics of these two parity cases probed by the double polarization asymmetry.

The dominance of  $C_1 C_3^*$  near threshold implies that the BT asymmetry  $\mathcal{D}_{xz}$  is not sensitive to the magnetic moment of the  $\Theta^+$ , since the  $u$ -channel contribution from the anomalous magnetic moment of  $\Theta^+$  to  $C_3$  is much smaller than the  $t$ -channel. We indeed see this by considering  $\kappa_\theta = \pm 1.91$ , which does not change the sign or basic features of  $\mathcal{D}_{xz}$ .

The dominance of  $C_1 C_3^*$  near threshold also holds even when  $K^*$  exchange included. As shown by the dashed and dotted curves in Fig. 3 at  $W = 2.1$  GeV,  $K^*$  exchange introduces significant interference into the asymmetries, which does not, however, change the dominant  $\sin \theta_{c.m.}$  behaviour. The explicit different signs of the BT asymmetry in the two cases, clearly distinguishes  $1/2^-$  from  $1/2^+$  parities near threshold.

Such a feature is undoubtedly remarkable, and experimentally sensible. Recall that a strong  $K^*$  coupling could lead to a significant enhancement of the cross sections [43, 44, 45, 46, 54] for the production of  $1/2^-$ . As a result, a single measurement of the cross section may not be sufficient for the determination of the  $\Theta^+$ 's parity. It turns out that the BT polarization asymmetry  $\mathcal{D}_{xz}$  is able to pick out the most distinct dynamical differences between  $1/2^+$  and  $1/2^-$ , and distinguish these two parities in a measurement near threshold.

It is worth noting that a similar behaviour of the Born terms also exists in the polarized beam asymmetry between these two parities [54]. However, the differences picked out by this single polarization asymmetry turn out not to be as marked as in the double polarization. Near threshold, although the Born terms will lead to differently signed asymmetries for  $1/2^+$  and  $1/2^-$  [54], they will destructively interfere with other mechanisms, such as the  $K^*$  exchange, and lose their characteristic feature.

In summary, we have analyzed the double polarization asymmetry,  $\mathcal{D}_{xz}$ , in  $\gamma n \rightarrow \Theta^+ K^-$ , and showed it to be a useful filter for determining the parity of  $\Theta^+$ , provided its spin-parity is either  $1/2^+$  or  $1/2^-$ . Due to dynamical reasons, asymmetry  $\mathcal{D}_{xz}$  near threshold would exhibit a similar behaviour but opposite sign. The advantage of studying polarization observables is that uncertainties arising from the unknown form factors can be partially avoided in a phenomenology. Therefore, although better knowledge of the form factors will improve the quantitative predictions, it should not change the threshold behaviour of  $\mathcal{D}_{xz}$  dramatically. However, special caution should be given to the roles played by a possible spin-3/2 partner in the  $u$ -channel, as well as  $s$ -channel nucleon resonances. In particular, as studied by Dudek and Close [41], the spin-3/2 partner may have a mass close to the  $\Theta^+$ . Thus, a significant contribution from the spin-3/2 pentaquark state may be possible. Its impact on the BT asymmetry needs to be investigated. In brief, due to the lack of knowledge in this area, any results for the BT asymmetry would be extremely important for progress in gaining insights into the nature of pentaquark states and dynamics for their productions. Experimental facilities at Spring-8, JLab, ELSA, and ESRF should have access to the BT asymmetry observable.

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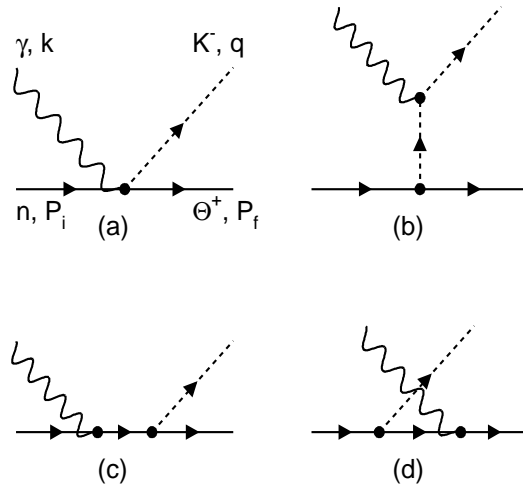


FIG. 1: Feynman diagrams for  $\Theta^+$  photoproduction in the Born approximation.

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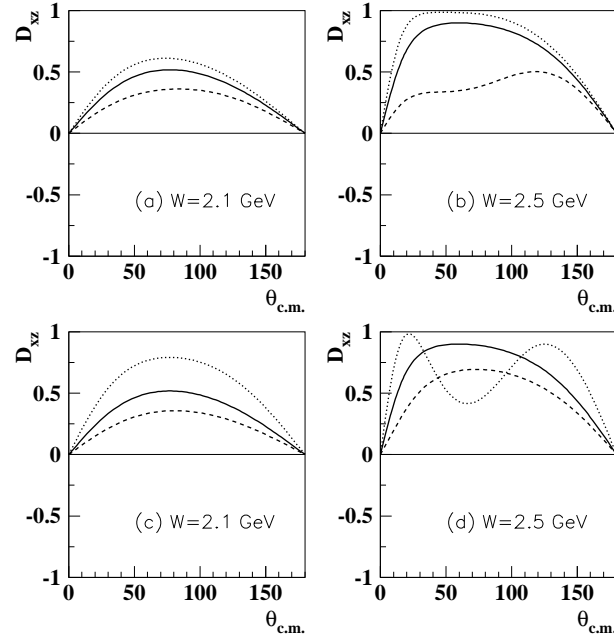


FIG. 2: Beam-target double polarization asymmetry for  $\Theta^+$  of  $1/2^+$  at  $W = 2.1$  and  $2.5$  GeV. The solid curves are results in the Born limit, while the dashed and dotted curves denote results with the  $K^*$  exchange included with different phases:  $(g_{\Theta NK^*}, \kappa_{\theta}^*) = (-2.8, -3.71)$  (dashed curves in (a) and (b)),  $(+2.8, +3.71)$  (dotted curves in (a) and (b)),  $(-2.8, +3.71)$  (dashed curves in (c) and (d)), and  $(+2.8, -3.71)$  (dotted curves in (c) and (d)).



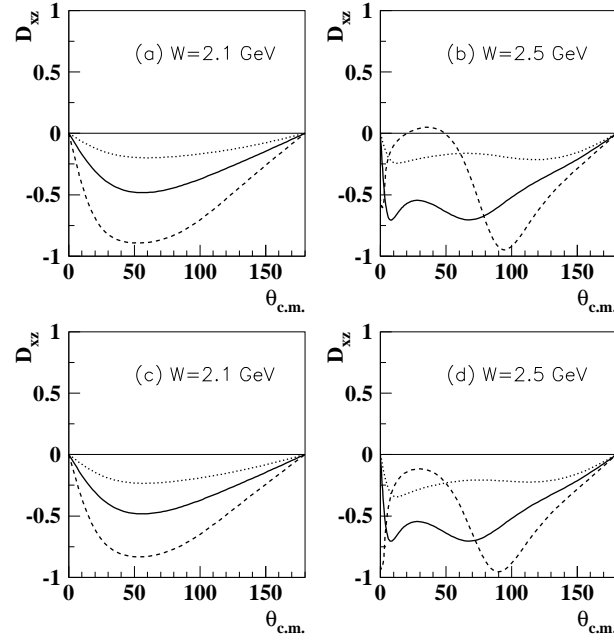


FIG. 3: Beam-target double polarization asymmetry for  $\Theta^+$  of  $1/2^-$  at  $W = 2.1$  and  $2.5$  GeV. The solid curves are results in the Born limit, while the dashed and dotted curves denote results with the  $K^*$  exchange included with different phases:  $(g_{\Theta NK^*}, \kappa_{\theta}^*) = (-0.61, -0.371)$  (dashed curves in (a) and (b)),  $(+0.61, +0.371)$  (dotted curves in (a) and (b)),  $(-0.61, +0.371)$  (dashed curves in (c) and (d)), and  $(+0.61, -0.371)$  (dotted curves in (c) and (d)).